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Particle Dark Matter*

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Abstract

There is plenty of evidence that most matter in the universe is dark (non–luminous). Particle physics offers several possible explanations. In this talk I focus on cold dark matter; the most promising candidates are then axions and the lightest supersymmetric particle. I briefly summarize estimates for the present relic density of these particles, and describe efforts to detect them.

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1) Introduction

"Dark Matter" (DM) is matter that does not emit detectable quantities of electromagnetic radiation; it nevertheless manifests itself through its gravitational pull. Historical examples are the planets Uranus and Pluto prior to their discovery by optical telescopes. In this talk I am only concerned with DM that covers volumes of galactic size or larger. Evidence for this kind of DM was first collected in the 1920's and 30's, but its existence became widely accepted only in the 1970's [1]. The least controversial evidence comes from "galactic rotation curves". Here one measures the rotational velocity of globular clusters, hydrogen clouds, or other objects, around spiral galaxies. Assuming that these objects are in stable orbits around their parent galaxies, Kepler's law tells us that the rotational velocity $v_{\rm rot}$ should decrease $\propto 1/\sqrt{r}$ at large distance r to the center of the galaxy, *if* the mass of the galaxy is concetrated in its visible part. However, observationally all rotation curves become essentially independent of r at large r, out to the largest observable distances; this implies that the mass inside the radius r grows linearly with r, i.e. the mass density drops $\propto 1/r^2$. Cosmologists like to express the average mass density of the universe in units of the critical or closure density $\rho_c \equiv (3H^2)/(8\pi G_N)$, where H is the Hubble constant and G_N is Newton's constant. Numerically, $\rho_c \simeq 2 \cdot 10^{-29} h^2$ g/cm^3 $\simeq 1.1 \cdot 10^{-5} h^2$ GeV/cm^3, where $h \equiv H/(100$ km/sec·Mpc). In these units, galactic DM halos contribute at least $\Omega \equiv \rho/\rho_c = 0.05$ to 0.1. Note that this is a lower bound, since we do not know where these halos end. Indeed, there is considerable evidence for larger values of Ω, extending to $\Omega = 1$, from studies of clusters of galaxies, the "streaming" of large numbers of galaxies, etc [2].

Luminous matter (stars, gas, dust) only contributes $\Omega = 0.01$ or slightly less. Further, Big Bang nucleosynthesis can only explain the observed abundances of light isotopes (D, ^{3}He, ^{4}He, ^{7}Li) if $0.01 \leq \Omega_b h^2 \leq 0.015$, where Ω_b is the *total* baryonic mass density in the universe. Direct measurements of the Hubble constant give $0.4 \leq h < 1$. Hence most DM must be non–baryonic, especially if the total $\Omega = 1$, as favoured by naturalness arguments and predicted by most inflationary models [3]. Note, however, that this range for Ω_b also implies the existence of baryonic DM, unless h is very close to 1. In some sense the recent observation of MACHOs [4] therefore strengthens the argument for non–baryonic DM: MACHOs are a good candidate for the predicted baryonic matter; having made a successful prediction obviously makes the overall model that much more trustworthy.

In order to be able to estimate detection rates of particle DM, we need to know the *local* DM flux, i.e. the local density and velocity of DM objects. Their density can be estimated by feeding various observations about our galaxy, including the observed rate of microlensing (MACHO) events, into a galactic model. A recent study by Turner et al. [5] quotes 0.2 GeV/cm^3 $\leq \rho_{\rm DM}^{\rm local} \leq$0.5 GeV/cm^3 (for a flattened halo, which is the preferred solution) for the non–MACHO DM density, although an all–MACHO halo cannot be excluded completely. Note that this is at least four orders of magnitude larger than the universal DM density (for $\Omega = 1$).

While the local DM density can at least be constrained from direct observations, estimates of the velocity distribution of DM particles are based almost entirely on galactic modelling. The simplest ansatz for the DM halo is an isothermal sphere. This leads to a Maxwellian velocity distribution in the glactic rest frame, with velocity dispersion $\langle v_{\rm DM}^2 \rangle^{0.5} \simeq 270$ km/sec $\simeq 10^{-3} c$. Note that the solar system moves through this isotropic DM soup with a velocity of about 220 km/sec. The DM velocity distribution on Earth is therefore highly anisotropic [6]: Most DM particles should come from the direction in which the solar system is moving. In

addition, the Earth moves around the Sun at about 15 km/sec; this leads to a small annual modulation of the DM velocity distribution as seen on Earth. Both the overall anisotropy and the annual modulation can in principle be used to suppress backgrounds in certain direct DM searches (see sec. 3).† Together with the value for $\rho_{\rm DM}^{\rm local}$ given in the previous paragraph, this gives a local DM flux $\Phi_{\rm DM} \simeq 10^5/({\rm cm}^2{\rm sec}) \cdot 100$ GeV$/m_{\rm DM}$, where $m_{\rm DM}$ is the mass of the DM particle. For comparison, the flux of cosmic ray muons at sea level is about $1/({\rm cm}^2{\rm sec})$; however, DM particles are much more difficult to detect than muons.

I focus here on "cold" DM (CDM), which was already non–relativistic when structure formation in the universe began. Scenarios with exclusively or dominantly "hot" (relativistic) DM are essentially excluded [9], unless the "seed" for structure formation comes from exotic objects like cosmic strings, rather than from quantum fluctuations during inflation, as commonly assumed.‡

Particle physics offers two main classes of CDM candidates, axions (sec. 2) and weakly interacting massive particles or WIMPs (sec. 3).

2) Axions

Let me start by briefly describing axion dark matter. Axions are the Goldstone bosons of the hypothetical global "Peccei–Quinn" (PQ) symmetry, which has originally been invented to solve the strong CP problem [11]. This new $U(1)$ symmetry allows to rotate the CP–violating phase θ away. The PQ symmetry is broken spontaneoulsy at scale f_a, and explicitly by QCD condensates that break chiral symmetry. This latter contribution gives the axion a finite mass [12]:

$$m_a \simeq 6.3 \text{ meV} \cdot \frac{10^9 \text{ GeV}}{f_a}. \tag{1}$$

The couplings of the axion to ordinary matter also scale like the inverse of f_a. Negative laboratory searches can be translated into upper bounds on these couplings, i.e. lower bounds on f_a. These laboratory bounds require the axion couplings to be so small that an axion, once produced, can transverse an entire star without interacting. Axion emission can therefore lead to too rapid cooling of stars unless the axion production rate is sufficiently small. The best bounds come from red giants and supernovae; a conservative limit is [12]

$$f_a > 10^9 \text{ GeV} \quad ==> \quad m_a \leq 6 \text{ meV}. \tag{2}$$

An upper bound on f_a can be derived from the requirement that axions do not overclose the universe. The derivation of this bound is not entirely straightforward since axions, unlike the WIMPs discussed in the next section, are not thermal relics. Rather, the main source of axions are either coherent field oscillations, or axion emission from global (axionic) strings that are produced during the PQ phase transition. The first of these sources leads to the bound [13]

$$f_a \leq h^2 \cdot 10^{12\pm 0.5} \text{ GeV}/\langle \theta_a^2 \rangle, \tag{3}$$

†Recently Cowsik et al. [7] suggested that the velocity dispersion is at least 600 km/sec. However, their analysis has been criticized by Gates et al. [8]. I will therefore stick to the "canonical" value for the time being.

‡There is some evidence that some 20% of all DM might be hot, e.g. neutrinos with masses in the eV range [9]. However, other models with only CDM work just as well [10].

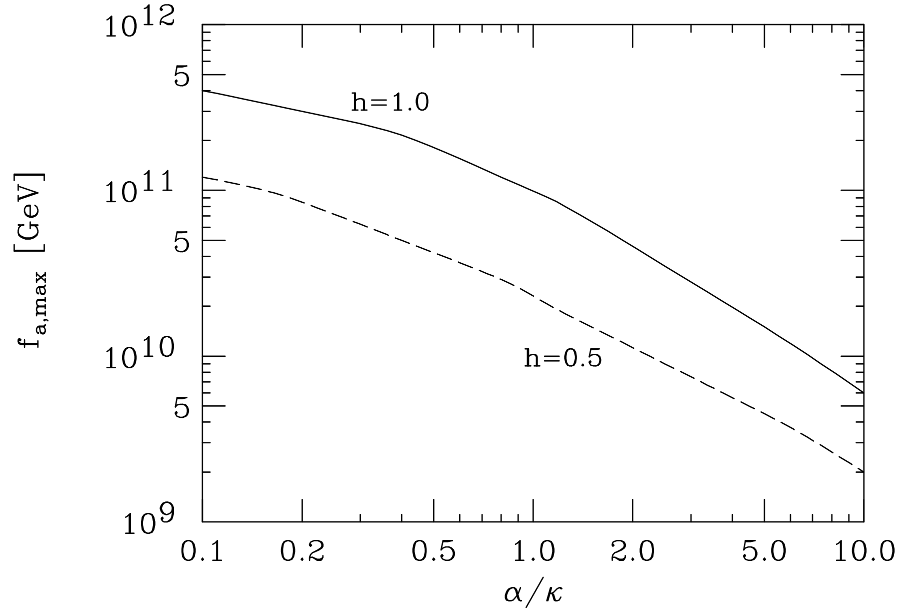

Figure 1: The upper bound on f_a is plotted as a function of the ratio α/κ explained in the text, for two values of the rescaled Hubble constant. It has been assumed that the present axion density is dominated by axion radiation off cosmic strings. Adapted from ref.[13].

where h is again the scaled Hubble constant, the uncertainty in the exponent is due to our limited understanding of details of the QCD phase transition, and the squared "vacuum mis-alignment angle" $\langle \theta_a^2 \rangle$ describes how far the axion field was from the origin of the potential at the onset of the QCD phase transition. Notice that the bound (3) becomes very weak if for some reason $\langle \theta_a^2 \rangle \ll 1$.

The PQ phase transition produces a network of global (axionic) strings. They are created in a rather energetic state, and lose energy by emitting axions. Short string loops are found to be the most potent source of axions. The upper bound on f_a that results from this source of axions is depicted in Fig. 1, taken from ref.[13]. Here the bound is plotted as a function of the ratio α/κ, where α describes the typical length of string loops, and the "backreaction parameter" κ describes the radiation power from strings per unit length; one expects $0.1 \leq \alpha/\kappa \leq 1$. The bound of Fig. 1 is clearly stronger than that of eq.(3); however, it can be evaded completely in inflationary models, if the reheating temperature is less than f_a. In this case no new strings are produced after inflation; if any strings were produced before inflation, their density would be so diluted that they contribute negligibly to total axion production. In such models the bound (3) would therefore be the relevant one.

However, the derivation of this bound assumes standard evolution of the Universe between the QCD phase transition and the present time. It could be weakened considerably if some mechanism produced a lot of additional entropy in that period. The reason is that one actually computes the axion to photon ratio just after the QCD phase transition, and then scales the photon density to the present one. Clearly this ratio would be diluted, i.e. the bound on f_a would be weakened, if an additional source produces a large number of photons after the QCD phase transition, thereby also increasing the entropy of the Universe. This is possible in supersymmetric axion models, where the late decay of the axion's fermionic partner, the "saxino", can weaken the bound on f_a by as much as a factor 10^3 [14]. Depending on assumptions, the upper bound on f_a is therefore somewhere in between a few times 10^{10} and

$\sim 10^{15}$ GeV; equivalently, m_a has to be in the range

$$6 \text{ meV} \geq m_a \geq (10^{-2} \text{ to } 10^2)\ \mu\text{eV}, \tag{4}$$

with the narrower window (stronger lower bound) corresponding to more standard scenarios.

Particles of such a small mass and non–relativistic velocity (see. Sec.1) cannot be detected by scattering. Rather, one searches for axion $\rightarrow$ photon conversion in a strong magnetic field [15]. This conversion can occur because 1–loop corrections produce an $a\gamma\gamma$ coupling of the form $g_{a\gamma\gamma} a F^{\mu\nu} \widetilde{F}_{\mu\nu}$. The strength of this coupling is quite model dependent, since even superheavy particles (with mass $\sim f_a$) can contribute in the loop. Typically,

$$g_{a\gamma\gamma} = g_\gamma \frac{\alpha_{\text{em}}}{\pi f_a}, \qquad \text{with } \frac{1}{3} \leq |g_\gamma| \leq 1. \tag{5}$$

This technique is now being pursued [16] by an experiment at Lawrence Livermore National Lab, for axion masses in the μeV range. Axions of this mass will convert into microwave photons, which can be detected only at very low temperatures. Further, a detectable level of microwave power will only be produced if the conversion is resonant, i.e. occurs in a cavity whose resonance frequency satisfies $\omega_{\text{res}} = m_a c^2/\hbar$. In order to scan a range of axion masses one therefore has to be able to change ω_{res} continuously, e.g. by using tuning rods. The strength of the axion signal (microwave power P_a) is then proportional to [16]

$$P_a \propto g_\gamma^2 V B_0^2 \frac{\rho_a}{m_a} Q, \tag{6}$$

where V is the volume of the cavity, B_0 is the strength of the applied $B-$field, ρ_a/m_a is the ambient axion number density, and Q characterizes the quality of the cavity (it is inversely proportional to the width of the cavity's resonance). The experiment now being run at LLNL [16] uses $B_0 = 8.5$ T, Q = (a few) 10^5, and $V \simeq 0.25$ m^3. It is expected to probe $|g_\gamma| \geq 1$ for 1.3 μeV $\leq m_a \leq$ 13 μeV. Although several orders of magnitude more sensitive than earlier pilot experiments, its sensitivity in terms of both g_γ and m_a is still marginal, see eqs.(4) and (5). In particular, the covered range of m_a is already excluded in "standard" axion cosmology [13], which corresponds to the more stringent lower bound in (3).

Finally, without going into detail I mention that axions occur quite naturally in superstring theory [17].

3) WIMPs

Unlike axions, weakly interacting massive particles (WIMPs) are thermal relics. Shortly after the Big Bang they were in thermal equilibrium with the soup of SM particles, i.e. the WIMP density was essentially given by the Maxwell–Boltzmann distribution. However, the rate for reactions that convert WIMPs into SM particles or vice versa drops quickly once the temperature of the Universe is less than the WIMP mass m_χ. Eventually this reaction rate will become smaller than the expansion rate of the Universe, at which point the WIMPs "freeze out", i.e. their density per co–moving volume remains essentially constant.

In order to compute the relic WIMP density $\Omega_\chi h^2$ one first introduces the rescaled inverse freeze–out temperature $x_f \equiv m_\chi/T_f$; it satisfies [3]

$$x_f = \ln \frac{0.038 g_{\text{eff}} M_p m_\chi \langle \sigma_{\text{eff}} v \rangle (x_f)}{\sqrt{g_* x_f}}. \tag{7}$$

Here, $g_{\rm eff}$ is the effective number of WIMP degrees of freedom (e.g., 2 for a single Majorana fermion), $M_P = 1.22 \cdot 10^{19}$ GeV is the Planck mass, and g_* is the effective number of relativistic degrees of freedom at temperature T_f; typically $g_* \simeq 80$ or so. Finally, the thermal average over the product of the effective WIMP annihilation cross section $\sigma_{\rm eff} = \sigma(\chi\chi \rightarrow \text{anything})$ and the relative velocity v between the two annihilating WIMPs is given by [3]

$$\langle \sigma_{\rm eff} v \rangle (x) = \frac{x^{1.5}}{2\sqrt{\pi}} \int_0^\infty dv v^2 e^{-v^2 x/4} \sigma_{\rm eff}(v) \cdot v, \tag{8}$$

where I have made use of the fact that WIMPs are non–relativistic at freeze–out ($x_f \simeq 20$). Eqs.(7) and (8) are usually solved by numerical iteration. Once x_f has been determined, the relic density is given by

$$\Omega_\chi h^2 = \frac{1.07 \cdot 10^9 \text{ GeV}^{-1}}{J(x_f)\sqrt{g_*} M_P}, \tag{9}$$

where the "annihilation integral" J is given by

$$J(x_f) = \int_{x_f}^\infty \frac{\langle \sigma_{\rm eff} v \rangle (x)}{x^2} dx. \tag{10}$$

In many cases it is sufficient to use a non–relativistic expansion of the annihilation cross section:

$$\sigma_{\rm eff} \cdot v = a + b v^2 + \cdots . \tag{11}$$

The integrations in eqs.(8) and (10) can then be performed analytically, with the result [3]

$$\langle \sigma_{\rm eff} v \rangle (x_f) = a + \frac{6b}{x_f} + \cdots ; \tag{12a}$$

$$J(x_f) = \frac{a}{x_f} + \frac{3b}{x_f^2} + \cdots . \tag{12b}$$

However, this approximation breaks down if $\sigma_{\rm eff}$ depends sensitively on v, e.g. in the vicinity of $s-$channel poles or near a threshold. Further, in some cases the WIMP stays in *relative* thermal equilibrium with another, slightly heavier new particle χ' even after it has dropped out of equilibrium with SM particles. In this case "co–annihilation" reactions of the form $\chi\chi' \rightarrow$ (SM particles) have to be included into the calculation of $\sigma_{\rm eff}$. I refer the reader to ref.[18] for a discussion of these more complicated cases.

Eq.(7) shows that the freeze–out temperature depends only logarithmically on the annihilation cross section. All potentially realistic WIMP candidates therefore have $x_f \simeq 20$, which corresponds to $v \simeq 1/3$, or $J(x_f) \simeq \sigma_{\rm eff} v/20$ if the expansion (11) can be used. Plugging this into eq.(9) one finds that $\Omega_\chi h^2$ is of order unity, i.e. χ makes a good CDM candidate, if

$$\sigma_{\rm eff} \simeq 2 \cdot 10^{-10} \text{ GeV}^{-2} \sim \mathcal{O}(0.1) \text{ pb}, \tag{13}$$

which is similar to typical weak interaction cross sections! This is a highly nontrivial "coincidence"; notice, e.g., the factors of M_P in eqs.(7) and (9), which somehow conspire with other numerical factors to single out the weak scale. It means that we do not have to invent either very strong or very weak new interactions in order to explain the existence of Dark Matter; any stable particle with "ordinary" weak interactions will do. (Hence the WI in WIMP, of course.)

A WIMP that can annihilate into SM particles with roughly weak interaction strength usually also scatters off ordinary matter with roughly weak interaction cross sections. The two most promising WIMP search techniques make use of these interactions.

"Direct" WIMP searches look for energy deposited by a WIMP inside a detector in a laboratory (situated underground to reduce cosmic ray induced backgrounds). Given that ambient CDM particles are expected to have velocity $v_\chi \simeq 10^{-3}c$, see Sec. 1, the deposited energy can only be of order $E_{\rm vis} \leq 50 \text{ keV} \cdot m_\chi/(100 \text{ GeV})$. Such a small energy deposition can be detected calorimetrically only if the detector is cooled. This technique was employed by the first experiments, using Germanium detectors that were originally designed to search for neutrino–less double β decay [19]. However, it is difficult to use this technique for detectors weighing more than a pound or so. More recently, people have therefore started to employ scintillating detectors, e.g. NaI. The currently most stringent direct detection limits come from such devices [20]; they restrict the WIMP scattering rate to be below a few tens of events/(kg·day). Future improvements are expected when the detectors are cooled to liquid nitrogen temperature. This should improve the discrimination between backgrounds due to β and γ radioactiv decays in or near the detector, and the signal due to nuclear recoil, which produces a slightly different light curve in the scintillator [21]. Other scintillating materials are also being explored; liquid Xenon appears to be particularly promising for large scale detectors [22].

Direct WIMP detection becomes difficult if m_χ greatly exceeds 100 GeV. On the one hand, the WIMP flux decreases like $1/m_\chi$ for given ambient CDM mass density. Further, for heavy WIMPs and not very small nuclei the momentum transfer can be large enough to lead to significant nuclear form factor suppression of the scattering cross section. Fortunately the second promising WIMP search strategy actually works better for heavier WIMPs.

This "indirect detection" technique is based on the observation [23] that WIMPs can be trapped inside the Sun or Earth, if they lose a sufficient amount of energy in a scattering reaction while travelling through these celestial bodies to become gravitationally bound to them. Such WIMPs will become concentrated near the center of these bodies. Eventually equilibrium will be reached between the rate of WIMP annihilation in and WIMP capture by these bodies, i.e. the annihilation rate becomes one half the capture rate (note that each annihilation destroys two WIMPs). Most annihilation products will be absorbed immediately in the surrounding medium. However, neutrinos can escape, and might then be detected in underground neutrino detectors on Earth. The by far largest signal comes from muon–neutrinos, since they can be converted into a muon well outside the detector proper, thereby greatly increasing the effective target volume. Furthermore, muons point back into the direction of the parent neutrino. One can therefore enhance the sensitivity of the search by comparing the neutrino flux from the center of the Sun or Earth with side bins. The expression for this indirect WIMP detection rate also receives a factor $1/m_\chi$ from the CDM flux. However, for $m_\chi < 1$ TeV this suppression is more than compensated by two enhancement factors: The $\nu_\mu \rightarrow \mu$ conversion cross section rises essentially linearly with energy (as long as $m_p m_\chi \ll m_W^2$, where m_p is the proton mass); and the range of muons produced in the rock (or ice) surrounding the detector, and hence the effective detector volume for ν_μ's coming from a fixed direction, increases linearly with $E_\mu \propto m_\chi$. Increasing m_χ beyond a TeV or so decreases the rate again, mostly because it becomes less likely for very heavy WIMPs to become trapped gravitationally after scattering off a nucleus in the Sun or Earth.

Both the direct and the indirect WIMP detection rate scale linearly with the WIMP–matter

scattering cross section.§ Since ambient CDM particles are non–relativistic, this cross section grows with the available center–of–mass energy $\sqrt{s}$. Scattering off electrons can therefore almost always be neglected compared to scattering off nuclei. One then distinguishes between spin–dependent interactions, through a pseudoscalar or axial vector coupling to the spin of the nucleus in question, and spin–indenpedent scalar or vector interactions. The latter can couple coherently to an entire nucleus, leading to an extra enhancement factor A^2 compared to the former, where A is the nucleon number of the target nucleus. As already noted above, if both the WIMP and the nucleus are heavy, nuclear form factor suppression has usually to be taken into account [6]. The best bounds on WIMPs with dominantly spin–dependent interactions comes from the NaI experiment mentioned above [20]. The best current bounds on WIMPs with mostly spin–independent couplings come from indirect searches at the underground detectors Kamiokande [24] and Baksan [25].

The perhaps most obvious WIMP candidate is a heavy neutrino. However, this suffers from both theoretical and experimental problems. Theoretically, it is not at all clear why a heavy neutrino should be stable, given that all known heavy SM fermions decay very rapidly. Further, an $SU(2)$ doublet (Dirac or Majorana) neutrino would have been found by existing searches, if its mass exceeds about 10 GeV (which is required by LEP data), and if it constitutes a substantial fraction of the DM halo of our galaxy. Actually, such a doublet neutrino does *not* make a good DM candidate unless it is very heavy; its relic density is too low, due to strong annihilation into W and Z pairs, unless its mass is in the TeV range. One way out is to introduce an $SU(2)$ singlet fermion, which mixes with the doublet after $SU(2)$ is broken. If the lighter eigenstate is stable, one can arrange the mixing angle such that $\Omega_\chi h^2 \simeq 1$ as desired. Such models are not only contrived, they are also severly constrained by existing data [20, 24, 25].

The currently by far most popular WIMP candidate is therefore the lightest supersymmetric particle (LSP) [2]. If one only includes those terms in the Lagrangian that are necessary to produce quark and lepton masses, $R-$parity is conserved and the LSP is stable. It then has to be electrically neutral in order to avoid bounds from exotic isotope searches [26]. This leaves us with two candidates, a sneutrino and the lightest neutralino. Sneutrinos have similar annihilation cross sections and capture rates as Dirac neutrinos; they are therefore essentially excluded as WIMP candidates [27].

Neutralino dark matter has been studied quite extensively [2]. In the MSSM the neutralinos are mixtures of the bino, the neutral wino, and the two neutral higgsinos; extended models also include singlet higgsinos [28], but I will stick to the MSSM here. Under the usual assumption of gaugino mass unification the LSP can never be wino–like, since in this case the $SU(2)$ gaugino mass M_2 is about twice as large as the $U(1)$ gaugino mass M_1 at the weak scale. One thus has to consider three scenarios. If the higgsino mass parameter $|\mu| \gg M_1$, the LSP is mostly a gaugino (a photino if $M_1 \ll m_Z$, and mostly bino if $M_1 \geq m_Z$). In the opposite limit $M_1 \gg |\mu|$, the LSP is mostly higgsino. Finally, if $M_1 \simeq |\mu|$, the LSP is a strongly mixed state.

In order to decide whether the LSP makes a good CDM candidate one first has to estimate its relic density. The to date most complete list of LSP annihilation cross sections has been compiled in ref.[29]. Here all possible 2–body final states have been treated, at tree level and using the expansion of eqs.(11), (12). It was realized subsequently that this is not always

§In case of indirect detection this is true only if equilibrium between capture and annihilation has been reached. Prior to equilibrium the signal is exponentially suppressed and therefore essentially unobservable [23]; this is the case for the signal from the Earth's center for most WIMP candidates that are significantly heavier than iron nuclei.

sufficient for a reliable estimate of the relic density. First, if the LSP is higgsino–like, the lighter chargino and the next–to–lightest neutralino are also higgsino–like and only slightly heavier than the LSP. In this case co–anihilation between the LSP and these slightly heavier states can be important. This is true in particular for $|\mu| \leq m_W$. The reason is that the $Z\widetilde{\chi}_1^0\widetilde{\chi}_1^0$ coupling is very small for almost pure higgsinos, while the $Z\widetilde{\chi}_2^0\widetilde{\chi}_1^0$ and $W\widetilde{\chi}_1^\pm\widetilde{\chi}_1^0$ couplings have essentially full gauge strength. As a result, light higgsinos become uninteresting as CDM candidates [30].

Furthermore, $s-$channel exchange diagrams involving Higgs bosons are often quite important. In this case the non–relativistic expansion (11) becomes unreliable. In this expansion one assumes that LSPs annihilate essentially at rest, i.e. the center–of–mass energy $\sqrt{s} = 2m_\chi$. However, the thermal energy of the LSPs can often push $\sqrt{s}$ up to the mass m_H of the exchanged Higgs boson if $2m_\chi$ is not too much smaller than m_H, thereby greatly reducing $\Omega_\chi h^2$. Not surprisingly, this effect is especially important in models where other ($t-$ and $u-$ channel) contributions to the annihilation cross section are small [31].

The dashed curves in Figs. 2 and 3 show contours of constant $\Omega_\chi h^2$ including some co–annihilation effects, and using a careful treatment of $s-$channel exchange contributions following ref.[32]. For these plots I have assumed gaugino mass unification, and have used a common mass m_0, as well as a common $A-$parameter A_0, for all sfermions at the GUT scale $M_X = 2 \cdot 10^{16}$ GeV. At the weak scale squarks are then heavier than sleptons (by as much as a factor 3–5), and $SU(2)$ doublet sleptons are heavier than $SU(2)$ singlets. However, I have left the SUSY breaking contributions to the Higgs boson masses free. This allows me to treat the mass m_A of the pseudoscalar Higgs boson as well as μ as free parameters. I have taken $\tan\beta = 2$, $m_t = 170$ GeV, $A_0 = 0$ and $m_A = 500$ GeV; in Fig. 2, $m_0 = 200$ GeV, while Fig. 3 is for $m_0 = 400$ GeV.

In both figures the lower–left region of the (μ, M_2) plane is excluded by LEP searches for charginos and neutralinos. The lower–right corners are excluded since here the lighter stop eigenstate $\tilde{t}_1$ would be lighter than the lightest neutralino; note that the off–diagonal entries of the stop mass matrix grow $\propto \mu$. The region excluded by this constraint is larger in Fig. 2 since it has smaller diagonal stop masses (smaller m_0). Finally, in the top–right corner of Fig. 2, the lighter stau eigenstate $\widetilde{\tau}_1$ becomes lighter than the lightest neutralino. Note that the ratio of stop and neutralino masses increases with increasing M_2 since $m_{\tilde{t}_1}$ gets large positive contributions from gluino loops. In contrast, the stau to neutralino mass ratio decreases with increasing M_2, since $m_{\widetilde{\tau}_1}$ only receives small contributions from bino loops.

Figs. 2 and 3 show prominent effects from the H^0 and A exchange poles at $M_2 \simeq 500$ GeV, and from the merged Z and h^0 exchange poles at $M_2 \simeq 100$ GeV; here h^0 and H^0 are the light and heavy neutral Higgs scalars, respectively. For $m_0 = 200$ GeV, Fig. 2, only a small region of the plane (top right) is excluded by the requirement $\Omega_\chi h^2 \leq 1$. Increasing m_0 to 400 GeV, Fig. 3, enlarges this region dramatically. The reason is that gaugino–like LSPs annihilate dominantly through $t-$channel exchange of $SU(2)$ singlet sleptons, which have both the smallest mass and the largest hypercharge of all sfermions. In the gaugino region the relic density therefore depends sensitively on m_0, roughly $\propto m_0^4$ if $M_1^2 \ll m_0^2$. On the other hand, in the higgsino region, $|\mu| < M_1 \simeq M_2/2$, the relic density depends very little on m_0. For $|\mu| < m_W$ co–annihilation into fermion–antifermion pairs via W and Z exchange is the dominant process, while heavy higgsinos annihilate dominantly into W or Z pairs; neither of these reactions depends on sfermion masses.

However, the estimate of the relic density in the higgsino region could be changed significantly by two effects that are not taken into account in these figures. First, co–annihilation into

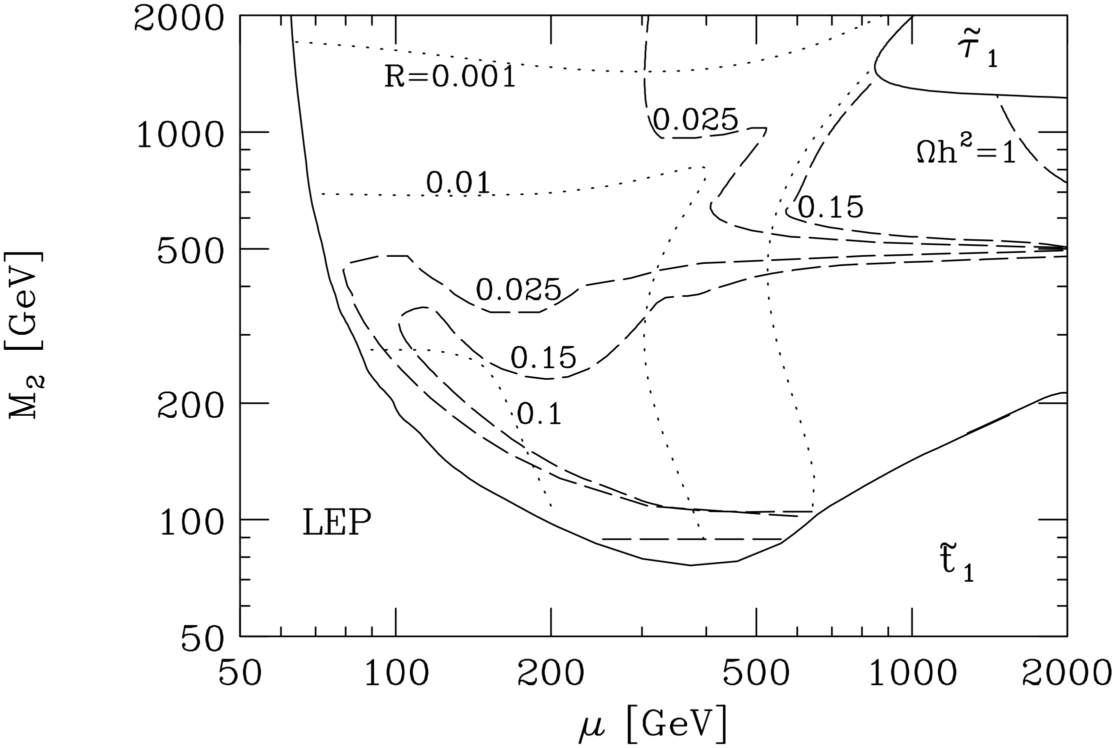

Figure 2: Contours of constant relic density (dashed), and of constant scattering rate in ^{76}Ge (dotted, in units of evts/(kg·day)). The values of the free parameters are: $m_t = 170$ GeV, tan$\beta = 2$, $A_0 = 0, m_A = 500$ GeV and $m_0 = 200$ GeV. The regions outside the solid curves are excluded, as described in the text.

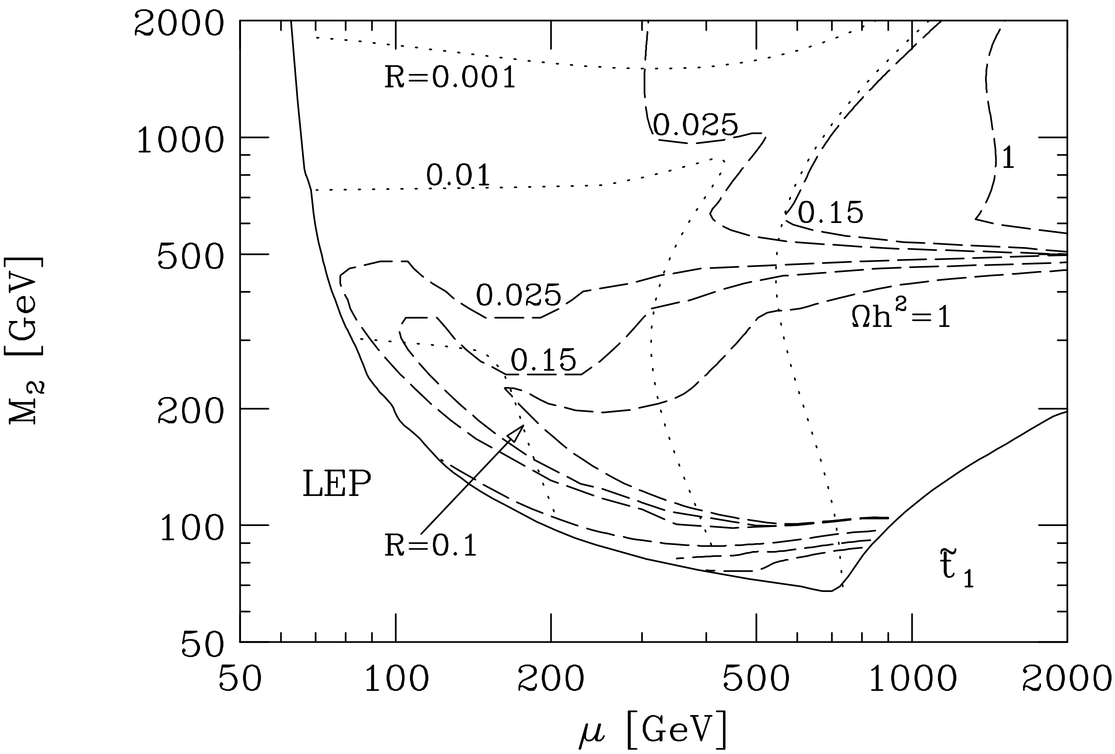

Figure 3: As Fig. 2, but for $m_0 = 400$ GeV.

states other than SM fermions or a photon and a W boson have not been included. Heavy higgsinos can co–annihilate into a host of other final states involving gauge or Higgs bosons; this will reduce their relic density even more. Secondly, it has recently been pointed out [33] that one–loop corrections can significantly change the mass splitting between higgsino–like states. These corrections can have either sign. If they increase the mass splitting, co–annihilation would be suppressed (its rate depends exponentially on the mass splitting), possibly reinstating light higgsinos as viable WIMP candidates.

Fig. 2 shows that a gaugino–like LSP usually does make a good CDM candidate, provided that m_0 is not too far from 200 GeV. In this context it is interesting to note that minimal supergravity models with heavy top quark favour a gaugino–like LSP, although a higgsino–like LSP remains possible even in these restrictive models. By "minimal supergravity" I mean models where the universality of soft breaking scalar masses is extended into the Higgs sector. The electroweak gauge symmetry is then more or less automatically broken by radiative corrections [34], which allows to determine $|\mu|$ and m_A in terms of m_t, tanβ, M_2 and A_0. One might hope that in these restrictive models the requirement $\Omega_\chi h^2 \leq 1$ allows to place upper limits on SUSY breaking parameters. Unfortunately this is not the case [29], as illustrated in Fig. 4. The solid curve in this plot shows $\Omega_\chi h^2$ as a function of tanβ for $m_0 = 500$ GeV, $A_0 = 0$, $m_t = 170$ GeV, $M_2 = 250$ GeV and $\mu < 0$. For small tanβ the relic density is indeed unacceptably high, due to the large value of m_0 chosen here. However, two effects reduce $\Omega_\chi h^2$ as tanβ is increased. First, the lighter $\tilde{\tau}$ eigenstate becomes lighter [35], due to loop corrections involving the τ Yukawa coupling (which grows with tanβ), and also due to increased $\tilde{\tau}_L - \tilde{\tau}_R$ mixing. Secondly, and even more importantly, m_A becomes smaller [35], due to loop effects involving the b and τ Yukawa couplings. For tan$\beta \simeq 48$, $2m_\chi \simeq m_A$ and the relic density is very small. One is thus forced to conclude that imposing an upper bound on the LSP relic density does not lead to strict upper bounds on sparticle masses, although it does significantly constrain the allowed parameter space.

Let me finally briefly discuss LSP detection. As mentioned earlier, both the direct and the indirect detection rate scale with the LSP–matter scattering cross section. The to date most comprehensive calculation of this cross section has been performed in ref.[36]. Since neutralinos are Majorana fermions, their vector current vanishes identically and their pseudoscalar current is proportional to the LSP velocity $\sim 10^{-3}$. This means that Z exchange only contributes to spin–dependent interactions, while pseudoscalar Higgs exchange can be neglected entirely. The exchange of scalar Higgs bosons gives rise to spin–independent interactions, and in fact often dominates the total scattering cross section. Finally, squark exchange contributes to both spin–dependent and spin–independent interactions. However, in models with gaugino mass unification and universal scalar masses at the GUT scale one has $m_{\tilde{q}} \geq 5 m_\chi$, in which case squark exchange contributions are quite small.¶

As an example, the dotted lines in Figs. 2 and 3 show contours of constant scattering rate in ^{76}Ge detectors in units of events/(kg·day), ignoring possible energy detection thresholds (which would reduce the rate), and assuming that the galactic halo is formed essentially entirely by LSPs. Recall that the present bounds on this rate lie above 10 evts/(kg·day). Clearly even increasing the sensitivity by a factor 100 only begins to scratch the depicted parameter space. Note also that in much of the region giving relatively large scattering rates for *fixed* local DM density the LSP does in fact not make a good CDM candidate, since $\Omega_\chi h^2$

¶The lighter stop eigenstate can be close to the LSP mass in such models, as noted above. However, there are no top quarks in the proton, so $\tilde{t}_1$ exchange can only contribute at 1–loop level, and remains quite small [36].

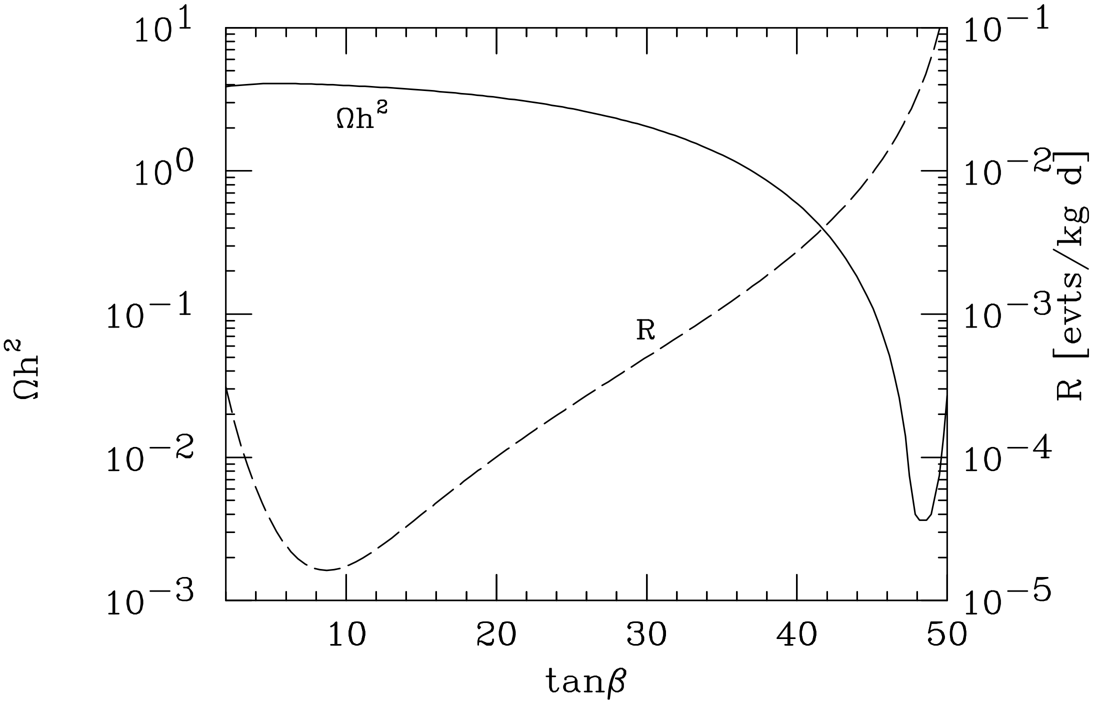

Figure 4: The LSP relic density (solid, referring to scale at left) and detection rate in ^{76}Ge (dashed, referring to scale at right) are plotted as a function of tanβ in a minimal supergravity model witgh radiative gauge symmetry breaking. The values of the other free parameters are $M_2 = 250$ GeV, $m_t = 170$ GeV, $A_0 = 0$ and $\mu < 0$.

is too low to even produce galactic haloes, which requires $\Omega_\chi h^2 \geq 0.025$ or so. The situation is even worse for $\mu < 0$, as shown by the dashed line in Fig. 4: Although $M_2 = 250$ GeV is quite moderate, the detection rate can be some five orders of magnitude below the present bound, due to destructive interference between h^0 and H^0 exchange. At large tanβ the rate increases, because the couplings of H^0 to d and s quarks grow $\propto$ tanβ, and also because $m_{H^0} \simeq M_A$ becomes smaller, as mentioned earlier. However, the mass of the charged Higgs boson is also reduced, which can lead to conflict with the upper bound on the branching ratio for $b \to s\gamma$ decays. Indeed, in ref.[37] it was pointed out that this last constraint makes it very difficult to construct viable SUSY models with large LSP detection rate even if all parameters are set "by hand" at the weak scale.[‖]

4) Summary and Conclusions

The argument for the existence of some exotic dark matter is quite compelling. Studies of structure formation in inflationary models indicate that most of this dark matter should have been "cold" already at the epoch of galaxy formation. Particle physics conveniently provides us with (at least) two CDM candidates: The axion (if f_a is chosen appropriately), and the supersymmetric LSP (if gaugino–like). Both have the attractive feature that they were originally introduced to solve a problem that has nothing to do with dark matter: Axions solve the strong CP problem, and supersymmetry solves the hierarchy problem. We have

‖The authors of ref.[38] did manage to do just that. However, they had to chose squark masses just above m_χ and m_A well below m_χ; neither assumption is particularly natural from a model building point of view.

seen that both candidates are quite difficult to detect (other than through their gravitational effects); experiments of the present and even next generations will only begin to probe the allowed parameter space. One should keep in mind that things could be even worse, though. For example, CDM could come from the "hidden sector" present in many supergravity or superstring models, in which case it would interact with normal matter *only* gravitationally,. Individual CDM particles would then be completely undetectable.

However, before we have to consider such depressing scenarios, we should fully explore the parameter space of the more attractive detectable candidates. In case of the LSP, a first decisive test will probably come from the LHC (or, if we are lucky, from LEP). However, even if the LHC finds a SUSY signal, detection of relic LSPs is necessary to prove that they are stable on cosmological time scales. (Collider experiments can only establish a lower bound of 10^{-8} seconds or so.) In contrast, I do not know of any realistic scheme to detect axions produced in the lab, given the extremely strong upper limits on their interaction strength; CDM axions are our only hope to detect these elusive particles, if they indeed exist. This once again illustrates the increasingly tight connection between particle physics and cosmology.

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